

2 Higgs or not 2 Higgs

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Motivated by recent results from the LHC experiments, we analyze Higgs couplings in two Higgs doublet models with an approximate PQ symmetry. Models of this kind can naturally accommodate sizable modifications to Higgs decay patterns while leaving production at hadron colliders untouched. Near the decoupling limit, we integrate out the heavy doublet to obtain the effective couplings of the SM-like Higgs and express these couplings in a physically transparent way, keeping all orders in (m_h/m_H) for small PQ breaking. Considering supersymmetric models, we show that the effects on the Higgs couplings are considerably constrained.

Introduction. It is with a light heart that we assume, as a working hypothesis, that the recent measurements presented by CMS and ATLAS [1, 2] hint to an Higgs boson of mass $m_h \sim 125$ GeV. Once this first step is taken, it is only human to collect σ/σ_{SM} best fit plots and feverishly inspect them for hints of new physics. Presently, the data are consistent with $\mathcal{O}(1)$ enhancements with respect to a Standard Model (SM) Higgs boson, for both gluon fusion (GF) and vector boson fusion (VBF) production channels, in the $\gamma\gamma$ [3, 4], and possibly also in the ZZ and WW decay modes [5, 6].

If the $\gamma\gamma$ rate (and potentially also ZZ and WW rates) turn out to be larger than in the SM, in both GF and VBF, then the multiple enhancements are more easily interpreted in terms of non-standard Higgs decay, rather than production. The simplest explanation being an $\mathcal{O}(1)$ reduction in the $h\bar{b}b$ coupling. However, a suppression in $h\bar{b}b$ must not be accompanied by a similar suppression in $h\bar{t}t$. Otherwise, without fortuitous interference between different new physics effects, GF production would be reduced by a similar amount. The situation is then such that (i) the couplings of h to down-type quarks and to up-type quarks exhibit non-universal sensitivity to new physics, and (ii) the effect in the Higgs-bottom coupling is more pronounced than in the Higgs-top coupling. These requirements are met in two Higgs-doublet models (2HDMs) with an approximate PQ symmetry.

In this class of models, some hierarchy, with one doublet parametrically heavier than the other, is motivated by experimental constraints. First, electroweak precision measurements limit the contribution of new physics to the ρ parameter. This implies that new physics at the TeV scale should approximately conserve the diagonal $SU(2)_c$ custodial symmetry. Since (H^-, A, H^+) transforms as a **3** of $SU(2)_c$, the splitting between m_{H^\pm} and m_A is constrained in these models [7, 8]. Furthermore the field $H \pm iA$ is charged under PQ , so the splitting between m_{H^\pm} and m_A is of the order of the PQ breaking, that we are assuming is moderate. Second, the mass of the charged Higgs boson H^\pm is constrained by its con-

tribution to the decay $B \rightarrow X_s \gamma$ and recent calculations give the bound [9]

$$m_{H^\pm} > 295 \text{ GeV} \quad (1)$$

at 95%CL. Without accidental cancellations, this bound also applies to models with a richer particle spectrum, such as supersymmetry. As a result, it is natural to expect the whole doublet $(H^+, H+iA)$ to be parametrically heavier than $m_h \sim 125$ GeV.

A mass hierarchy motivates an effective theory analysis of the 2HDM with only a single SM-like Higgs boson at the weak scale. In this paper we take on this analysis, adopting a somewhat different approach than in the previous literature. Focusing our attention to modified Higgs couplings, we present our results in a way that we find more transparently related to the symmetries of the theory from which the 2HDM and, eventually, the SM is assumed to descend. We show how these symmetries are reflected in observable Higgs signals and demonstrate the utility of our approach by easily deriving the modified Higgs couplings in several supersymmetric embeddings of the 2HDM.

2HDM analysis. Consider a type-II 2HDM with $H_{1,2} \sim (1, 2)_{+1/2}$, where only H_1 couples directly to $\bar{Q}_L d_R$ and only H_2 couples directly to $\bar{Q}_L u_R$ at high energies. Neglecting leptons for now, the Lagrangian is [10, 11]

$$\begin{aligned} -\mathcal{L} = & H_1^\dagger \mathcal{D}^2 H_1 + H_2^\dagger \mathcal{D}^2 H_2 + m_1^2 |H_1|^2 + m_2^2 |H_2|^2 \\ & + \frac{\lambda_1}{2} |H_1|^4 + \frac{\lambda_2}{2} |H_2|^4 + \lambda_3 |H_1|^2 |H_2|^2 + \lambda_4 |H_1 \sigma_2 H_2|^2 \\ & + \left\{ \frac{\lambda_5}{2} (H_1^\dagger H_2)^2 + (H_1^\dagger H_2) (m_{12}^2 + \lambda_6 |H_1|^2 + \lambda_7 |H_2|^2) \right. \\ & \left. + Y_u H_2 \bar{u}_R Q_L + Y_d H_1^\dagger \bar{d}_R Q_L + cc \right\}. \end{aligned} \quad (2)$$

The parameters m_{12}^2 , λ_6 , λ_7 and λ_5 violate a $U(1)_{PQ}$ under which $(H_1^\dagger H_2)$ has charge +1. A discrete Z_2 subgroup of this $U(1)_{PQ}$ controls the mixing between the two doublets. In this paper, we loosely refer to approximate Z_2 as the PQ limit. Since the coupling λ_5 is even under the

Z_2 , it does not need to be small for our analysis to apply and indeed we will treat it collectively with other Z_2 -even couplings. We parameterize spontaneous symmetry breaking (SSB) in a unitary gauge with

$$H_1 = \begin{pmatrix} h^+ \\ \frac{h_1 + ia}{\sqrt{2}} \end{pmatrix}, \quad H_2 = \begin{pmatrix} 0 \\ \frac{h_2}{\sqrt{2}} \end{pmatrix}, \quad \langle h_2 \rangle = v_2, \quad (3)$$

where a, h_1, h_2 and the VEV v_2 are real.

It is possible to diagonalize the Higgs mass matrix and express the couplings in terms of the rotation angle α connecting the interaction to the mass basis and of the ratio $\tan \beta$ between the VEVs of H_2 and H_1 . This procedure gives $r_d \equiv \frac{vg_{hd\bar{d}}}{m_d} = -(\sin \alpha / \cos \beta)$, $r_u \equiv \frac{vg_{hu\bar{u}}}{m_d} = (\cos \alpha / \sin \beta)$ and $r_V \equiv \frac{vg_{hVV}}{2m_V^2} = \sin(\beta - \alpha)$. The trigonometric expressions for the r_X 's are useful as they provide the exact result and make apparent simple algebraic relations between them [12]. They are less useful, however, if one looks for more insight into the underlying theory. Here, much in the spirit of [13], we abandon the exact but somewhat less revealing $\alpha - \beta$ formulation in favor of a perturbative expansion, keeping track of the couplings in Eq. (2) as we work out the solution.

Our scheme is useful if the doublet H_1 is heavier than H_2 , so that around the scale m_h only H_2 is accessible. With this framework in mind we will obtain an effective action for h_2 to order $(B/M_1^2)^3$, where

$$M_1^2 = m_1^2 + \frac{\lambda_{35}h_2^2}{2}, \quad B = m_{12}^2 + \frac{\lambda_7h_2^2}{2} \quad (4)$$

with¹ $\lambda_{35} = \lambda_3 + \lambda_5$. We will not need to assume that $\lambda_{35}v^2 \ll m_1^2$. This will improve the accuracy of our results for a mild hierarchy $m_h \lesssim m_H$.

Before proceeding to integrate out the heavy fields in H_1 , we note some simplifying properties of the Lagrangian. First, we assume that CP is conserved to a good approximation, and take all the potential couplings to be real. Under this assumption, scalars and pseudo scalars do not mix and we need only consider diagrams involving the two neutral scalars h_1 and h_2 . Second, as defined in Eq. (2), λ_4 projects neutral onto charged states and vice versa. It does not enter in tree diagrams with no charged external Higgs fields and we can ignore it in what follows. Third, working to $\mathcal{O}(B^3/M_1^6)$, we can ignore λ_6 and λ_1 that affect the results beginning at $\mathcal{O}(B^3/M_1^6)$ and $\mathcal{O}(B^4/M_1^8)$, respectively.

Integrating out h_1 we obtain

$$\begin{aligned} -\mathcal{L}_{eff} = & \frac{1}{2}h_2\mathcal{D}^2h_2 + \frac{1}{2}m_2^2h_2^2 + \frac{\lambda_2}{8}h_2^4 + \frac{Y_u}{\sqrt{2}}h_2t\bar{t} \\ & - \frac{1}{2}Bh_2\frac{1}{\mathcal{D}^2 + M_1^2}Bh_2 - \frac{Y_b}{\sqrt{2}}b\bar{b}\frac{1}{\mathcal{D}^2 + M_1^2}Bh_2. \end{aligned} \quad (5)$$

The interactions of the canonically normalized SM-like Higgs h with the fermions and gauge bosons can be read off from (5), after accounting for wave function renormalization at $\mathcal{O}(B^2/M_1^4)$. In particular, the bottom-Higgs Lagrangian is given by

$$\frac{Y_b}{\sqrt{2}}b\bar{b}\frac{1}{\square + M_1^2}B\left(v_2 + \left(1 - \frac{f'^2}{2}\right)h\right) \quad (6)$$

with

$$\begin{aligned} v_2 &= v\left(1 - \frac{f^2}{2v^2}\right), \quad v^2 = \frac{1}{\sqrt{2}G_F} \cong (246 \text{ GeV})^2, \\ f &= \left\langle \frac{Bh_2}{M_1^2} \right\rangle, \quad f' = \frac{\partial f}{\partial v_2}. \end{aligned} \quad (7)$$

Using (5) and (6) we obtain:

$$\begin{aligned} r_b &= \frac{vg_{hb\bar{b}}}{m_b} = \frac{r_t}{1 - \frac{m_h^2}{M_1^2}} \left(1 + \frac{\lambda_7v_2^2}{B} - \frac{\lambda_{35}v_2^2}{M_1^2 - m_h^2}\right), \\ r_t &= \frac{vg_{ht\bar{t}}}{m_t} = 1 + \frac{B^2}{2M_1^4} \left(1 - \left(1 + \frac{\lambda_7v_2^2}{B} - \frac{\lambda_{35}v_2^2}{M_1^2}\right)^2\right), \\ r_V &= \frac{vg_{hVV}}{2m_V^2} = 1 - \frac{B^2}{2M_1^4} \left(\frac{\lambda_7v_2^2}{B} - \frac{\lambda_{35}v_2^2}{M_1^2}\right)^2. \end{aligned} \quad (8)$$

The appearance of terms (m_h^2/M_1^2) in r_b is due to the derivative operator in the effective vertex (6). The \square operator is replaced by $\square \rightarrow -m_h^2$ when acting on an external Higgs particle and by $\square \rightarrow 0$ when acting on the vacuum. Since m_h corresponds to the physical mass, the \square operator automatically includes radiative corrections to m_h .

The deviation of r_t from unity is parametrically small, beginning at $\mathcal{O}(B^2/M_1^4)$. The deviation of r_V scales similarly. In contrast, the deviation of r_b does not scale with (B/M_1^2) . It can be parametrically $\mathcal{O}(1)$ provided that either (i) $\lambda_7v^2 \sim m_{12}^2$ or (ii) $\lambda_{35}v^2 \sim m_1^2$.

The condition $\lambda_7v^2 \sim m_{12}^2$ implies that the hard and soft breakings of the PQ are comparable at the scale of SSB. Note that it is perfectly possible to have $\lambda_7v^2 \sim m_{12}^2$ and $m_{12}^2 \ll m_1^2$. For instance, if the theory at some high scale has $m_{12}^2 \sim 0$ but finite λ_7 , we can expect $m_{12}^2 \sim \lambda_7m_1^2/(4\pi)^2$ at scales below m_1 . In this case we can have an $\mathcal{O}(1)$ correction to r_b while the heavier doublet can be very heavy, $m_H \sim \text{TeV}$. This shows that r_b is a sensitive probe for hard breaking of the PQ [13].

The second condition, $\lambda_{35}v^2 \sim m_1^2$, implies that a sizable part of the mass of H_1 is driven by SSB. This is the relevant condition for models in which hard breaking of the PQ is absent or small (like e.g. the MSSM). In this case, a discernible deviation of r_b from unity implies a light second doublet with $m_H \sim v$. The corrections $\sim (m_h^2/M_1^2)$ coming from the derivative expansion can then be relevant; note that these terms correct r_b with a definite positive sign. In the interesting case where soft

¹ Compared with the basis of [10], $(B/M_1^2) \sim 1/\tan \beta$ and our λ_{35} equals their λ_{345} .

PQ breaking is also small, $B \ll m_h^2$, we can expand r_b in (B/M_1^2) . In that case, we can replace $M_1^2 \rightarrow m_H^2$, valid to $\mathcal{O}(B^2/M_1^4)$, in Eq. (8), obtaining

$$r_b \approx \left(1 - \frac{m_h^2}{m_H^2}\right)^{-1} \left[1 - \left(1 - \frac{m_h^2}{m_H^2}\right)^{-1} \frac{\lambda_{35} v^2}{m_H^2}\right]. \quad (9)$$

Eq. (9) is correct to all orders in (m_h^2/m_H^2) and $(\lambda_{35} v^2/m_H^2)$ and to $\mathcal{O}(B^2/M_1^4)$.

Modifying the Higgs decay to bottom quarks affects the other search channels by changing the total width. A 125 GeV Higgs in the SM has $BR(h \rightarrow b\bar{b}) \approx 56\%$ so, for instance, the diphoton signal will be

$$\frac{\sigma \times BR(h \rightarrow \gamma\gamma)}{\sigma \times BR(h \rightarrow \gamma\gamma)_{\text{SM}}} \cong \frac{1}{1 + 0.56(r_b^2 - 1)}. \quad (10)$$

The effect on the ZZ, WW final states is similar. In Fig. 1 we plot the diphoton enhancement from Eqs. (9)-(10). The maximal enhancement is about a factor of two and is obtained for $m_H^2 = m_h^2 + \lambda_{35} v^2$. This means that, in the context of a 2HDM with an approximate PQ and for order one couplings $\lambda_{35} \sim 1$, taking the recent best fit ATLAS and CMS results [3, 4] at face value implies a light second doublet $m_H \sim 300$ GeV. Note that in Eq. (10) we neglected the charged Higgs loop contribution to the coupling $h\gamma\gamma$. In the appendix we show that this contribution is indeed negligible for the range of m_{H^\pm} that is consistent with $b \rightarrow X_s \gamma$.

Finally, using Eqs. (8) we give a quick prescription for computing the correction to r_b in models with a type-II 2HDM at low energies.

1. If the theory contains hard breaking of the PQ via λ_7 , then significant deviation is possible even for $m_H \sim \text{TeV}$ in which case the leading effect is [10, 13]

$$r_b \approx 1 + \frac{2}{1 + 2m_{12}^2/(\lambda_7 v^2)}. \quad (11)$$

2. If there is little or no hard breaking of the PQ , $\lambda_7 v^2 \ll m_{12}^2$, then a modified r_b requires a light second doublet. When soft PQ breaking is also small, $B \ll m_h^2$, Eq. (9) resums all powers of (m_h^2/m_H^2) and $(\lambda_{35} v^2/m_H^2)$.

So far we have neglected the Higgs coupling to leptons, but those can be added in a straightforward manner. If the doublet H_1 that couples to the down quarks couples also to the leptons, then $r_b = r_\tau$ and the change to the total width is amplified by a small factor $1 + (m_\tau/m_b)^2/3 \sim 1.1$.

Supersymmetric examples. We now examine supersymmetric extensions of the SM with a 2HDM effective theory near the weak scale and extract the modifications to Higgs observables.

In supersymmetry, holomorphy of the superpotential requires a second Higgs doublet in order to couple the

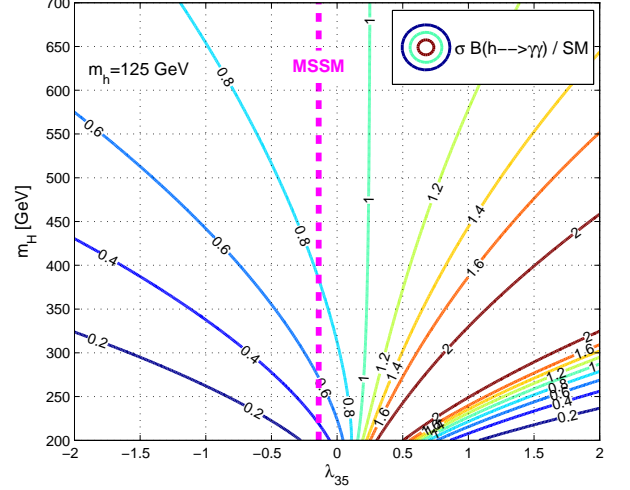


FIG. 1: Contours of $\sigma \times BR(h \rightarrow \gamma\gamma)/(SM)$ vs. λ_{35} and m_H , for $m_h = 125$ GeV and $\lambda_7 = 0$. (Recall $\lambda_{35} = \lambda_3 + \lambda_5$.) The MSSM prediction, neglecting loop corrections to the bottom Yukawa (but effectively including corrections to the Higgs potential), is shown by the dashed line.

Higgs sector to both up- and down-type quarks. Identifying $H_d = i\sigma_2 H_1^*$, $H_u = H_2$, the tree level quartic couplings of the MSSM are

$$\lambda_1 = \lambda_2 = \frac{g^2 + g'^2}{4}, \quad \lambda_3 = -\frac{g^2 + g'^2}{4}, \quad \lambda_4 = \frac{g^2}{2}, \quad \lambda_5 = \lambda_6 = \lambda_7 = 0. \quad (12)$$

The coupling $\lambda_{35} = \lambda_3 + \lambda_5 \approx -0.14$ is negative and so tends to increase $hb\bar{b}$. With λ_{35} fixed and assuming $m_h = 125$ GeV, Eq. (9) tells us that, neglecting loop corrections to the bottom Yukawa, the value of r_b depends only on m_H with little sensitivity to the details of the supersymmetric spectrum. This result is in good agreement with FeynHiggs [14]. The MSSM prediction for r_b in this case translates to the vertical dashed line in Fig. 1.

In some corners of the MSSM parameter space, our analysis ceases to give the leading result due to loop effects outside of the Higgs sector [10, 15–17]. The bottom Yukawa is corrected by stop-higgsino and sbottom-gaugino loops. The first contribution scales as $(\mu A_t/m_{soft}^2)/(4\pi)^2$, and can dominate the coupling for $(\mu A_t/m_{soft}^2) \sim \text{few}$. The second contribution scales as $(Y_b \mu m_\lambda/m_{soft}^2)/(4\pi)^2$ and can dominate for $(\mu m_\lambda/m_{soft}^2) \gtrsim (4\pi)^2 (B/m_H^2) \sim (4\pi)^2/\tan\beta$. In addition, light sfermions can affect the $h\gamma\gamma$ (and potentially hGG) vertex. In particular, very light staus could give an enhancement if (M_1^2/B) (or $\tan\beta$) is so large that the tau Yukawa becomes $\mathcal{O}(1)$ [15]. This defines a lower limit for (B/M_1^2) where we can neglect bottom and tau loop corrections: $(B/M_1^2) \gtrsim (\sqrt{2}m_b/v) \sim 1/40$.

It is interesting to ask whether simple extensions of the MSSM that accommodate a large Higgs mass in a more natural way can also reduce $hb\bar{b}$. We briefly examine three such models, the NMSSM, the BMSSM and a $U(1)_X$. Considering a Z_3 version of the NMSSM, we write the superpotential

$$\mathcal{W} = \lambda S H_u H_d + \frac{\kappa}{3} S^3. \quad (13)$$

We assume that S is somewhat heavy so that at low energy the theory can be described by the 2HDM. The coupling λ_3 is then given by

$$\lambda_3 = -\frac{g^2 + g'^2}{4} + |\lambda|^2 \approx -0.14 + 0.5 \left| \frac{\lambda}{0.7} \right|^2 \quad (14)$$

and is larger than in the MSSM. Still, as can be seen from Fig. 1, for this effect alone to achieve $r_b < 1$, $\lambda > 0.7$ is required, above the limit of perturbative unification [18].

Adding to Eq. (13) the supersymmetric mass terms $\mu H_u H_d$ and $\frac{\Lambda}{2} S^2$ with $\Lambda \gg \mu$ produces the BMSSM [7]. The spurion $\epsilon_1 = (\lambda^2 \mu^* / \Lambda)$ carries PQ charge -1 and induces $\lambda_6 = \lambda_7 = -2\epsilon_1$. For $(B/M_1^2) \lesssim 1/10$, ϵ_1 could decrease r_b while making a negligible correction to m_h [13].

Next, consider a gauge $U(1)_X$ extension under which the Higgs fields are charged, $q_{H_u} = q_{H_d} \equiv q_H$. The scalar potential receives a correction

$$V = V_{\text{MSSM}} + \frac{g_X^2 q_H^2}{2} (|H_2|^2 + |H_1|^2)^2 \quad (15)$$

Leading to $\lambda_3 \rightarrow \lambda_3 + g_X^2 q_H^2$. The modification to the $hb\bar{b}$ coupling in this example is in fact limited by the Higgs mass: since $\delta m_h^2 \sim 2g_X^2 q_H^2 v^2$, we should impose $g_X^2 q_H^2 < m_h^2 / (2v^2) \approx 0.12$. We therefore do not expect large deviations from the MSSM tree level predictions.

Conclusions. We have shown that a 2HDM close to the decoupling limit with an approximate PQ symmetry (or a Z_2 subgroup of it) can produce $\mathcal{O}(1)$ deviations in the lightest Higgs couplings to SM particles. Integrating out the heavier Higgs, we presented the effective couplings of the light SM-like Higgs in a physically transparent way. Keeping derivative operators in the expansion allowed us to include automatically radiative corrections to the lightest Higgs mass, important in the limit of small PQ breaking but a mild hierarchy $m_H \sim m_h$.

Our results are applicable to any type-II 2HDM not far from decoupling. In particular, considering the MSSM and some of its extensions, our analysis elucidates why it is hard to enhance $h \rightarrow \gamma\gamma$. If the current experimental hints are confirmed as the statistics increases, these results will allow to set bounds on the heavy Higgs mass m_H in large regions of these models parameter space.

Assuming a 2HDM with an approximate PQ and natural couplings, taking the recent best fit ATLAS and CMS results [3, 4] at face value implies a not too-heavy second doublet $m_H \sim 300$ GeV. Further support for this

possibility should come from better measurements of the SM-like Higgs decay patterns, where the generic type-II 2HDM predicts similar enhancements in gluon fusion and vector boson fusion production for $\gamma\gamma$, ZZ and WW final states, with correspondingly decreased $h \rightarrow b\bar{b}$.

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Appendix: Charged Higgs contribution to $h\gamma\gamma$. The decay $h \rightarrow \gamma\gamma$ is mediated by a dimension five Lagrangian, that we parametrize by [19]

$$\mathcal{L}_\gamma = -\frac{2\pi\alpha v c_\gamma}{\Lambda^2} h F_{\mu\nu} F^{\mu\nu} - \frac{2\pi\alpha v \tilde{c}_\gamma}{\Lambda^2} h \tilde{F}_{\mu\nu} F^{\mu\nu}. \quad (16)$$

In the absence of CP-violation, $\tilde{c}_\gamma = 0$. The contribution of the charged Higgs loop is given by [20]

$$\frac{c_\gamma}{\Lambda^2} = -\frac{\lambda_{hH^\pm H^\pm}}{v} \frac{f(\tau)}{24\pi^2 m_{H^\pm}^2}, \quad \tau = \frac{m_h^2}{4m_{H^\pm}^2} \quad (17)$$

with

$$f(r) = 3 \frac{\arcsin^2(\sqrt{r}) - r}{r^2} \cong 1 + 0.6r, \quad (18)$$

where the last approximation is valid to one percent accuracy for $m_h^2 / (4m_{H^\pm}^2) < 0.2$ or, in case of $m_h = 125$ GeV, for $m_{H^\pm} > 140$ GeV. The coupling $(\lambda_{hH^\pm H^\pm} / v)$ can be obtained from Eq. (2),

$$\frac{\lambda_{hH^\pm H^\pm}}{v} = \lambda_{34} \quad (19)$$

with $\lambda_{34} = \lambda_3 + \lambda_4$ to $\mathcal{O}(B^2/M_1^4)$. In the MSSM, for example, $\lambda_{34} = (g^2 - g'^2)/4 \approx 0.07$.

Adding this correction to the SM W and top loop contributions gives

$$\frac{\Gamma(h \rightarrow \gamma\gamma)}{\Gamma(h \rightarrow \gamma\gamma)_{\text{SM}}} = \left| \frac{r_V \mathcal{I}^W + r_t \frac{4}{3} \left(1 - \frac{\alpha_s}{\pi}\right) \mathcal{I}^t}{\mathcal{I}^\gamma} - \frac{4\pi^2 v^2 c_\gamma}{\Lambda^2 \mathcal{I}^\gamma} \right|^2, \quad (20)$$

with the loop function $\mathcal{I}^\gamma = \mathcal{I}^W + \frac{4}{3} \left(1 - \frac{\alpha_s}{\pi}\right) \mathcal{I}^t$ taken from [19]. For $m_h = 125$ GeV we find $\mathcal{I}^W = -2.09$, $\mathcal{I}_t = 0.34$, $\mathcal{I}^\gamma = -1.645$. Unless the charged Higgs is very light, or the relevant quartic couplings are large, the contribution to the $h\gamma\gamma$ coupling makes a negligibly small correction to the SM terms:

$$\frac{\Gamma(h \rightarrow \gamma\gamma)}{\Gamma(h \rightarrow \gamma\gamma)_{\text{SM}}} \cong \left| 1.27r_V - 0.27r_t - 0.05 \left(\frac{\lambda_{hH^\pm H^\pm}}{v} \right) \left(\frac{m_{H^\pm}}{350 \text{ GeV}} \right)^{-2} \right|^2.$$

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- [1] ATLAS Collaboration, arXiv:1202.1408 [hep-ex].
- [2] CMS Collaboration, arXiv:1202.1488 [hep-ex].
- [3] ATLAS Collaboration, arXiv:1202.1414 [hep-ex].
- [4] CMS Collaboration, arXiv:1202.1487 [hep-ex].
- [5] ATLAS Collaboration, arXiv:1202.1415 [hep-ex]. G. Aad *et al.* [ATLAS Collaboration], arXiv:1112.2577 [hep-ex].
- [6] CMS Collaboration, arXiv:1202.1489 [hep-ex]. CMS Collaboration, arXiv:1202.1416 [hep-ex].
- [7] M. Dine, N. Seiberg and S. Thomas, Phys. Rev. D **76**, 095004 (2007) [arXiv:0707.0005 [hep-ph]].
- [8] J. Mrazek, A. Pomarol, R. Rattazzi, M. Redi, J. Serra and A. Wulzer, Nucl. Phys. B **853**, 1 (2011) [arXiv:1105.5403 [hep-ph]].
- [9] M. Misiak, H. M. Asatrian, K. Bieri, M. Czakon, A. Czarnecki, T. Ewerth, A. Ferroglia and P. Gambino *et al.*, Phys. Rev. Lett. **98**, 022002 (2007) [hep-ph/0609232].
- [10] J. F. Gunion and H. E. Haber, Phys. Rev. D **67** (2003) 075019 [hep-ph/0207010].
- [11] G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, M. Sher and J. P. Silva, arXiv:1106.0034 [hep-ph].
- [12] P. M. Ferreira, R. Santos, M. Sher and J. P. Silva, arXiv:1112.3277 [hep-ph].
- [13] L. Randall, JHEP **0802**, 084 (2008) [arXiv:0711.4360 [hep-ph]].
- [14] S. Heinemeyer, W. Hollik and G. Weiglein, Comput. Phys. Commun. **124**, 76 (2000) [hep-ph/9812320]; S. Heinemeyer, W. Hollik and G. Weiglein, Eur. Phys. J. C **9**, 343 (1999) [hep-ph/9812472]; G. Degrandi, S. Heinemeyer, W. Hollik, P. Slavich and G. Weiglein, Eur. Phys. J. C **28**, 133 (2003) [hep-ph/0212020]; M. Frank, T. Hahn, S. Heinemeyer, W. Hollik, H. Rzehak and G. Weiglein, JHEP **0702**, 047 (2007) [hep-ph/0611326].
- [15] M. Carena, S. Gori, N. R. Shah and C. E. M. Wagner, arXiv:1112.3336 [hep-ph].
- [16] G. L. Kane, G. D. Kribs, S. P. Martin and J. D. Wells, Phys. Rev. D **53**, 213 (1996) [hep-ph/9508265].
- [17] M. S. Carena, J. R. Ellis, A. Pilaftsis and C. E. M. Wagner, Nucl. Phys. B **586**, 92 (2000) [hep-ph/0003180].
- [18] L. J. Hall, D. Pinner and J. T. Ruderman, arXiv:1112.2703 [hep-ph].
- [19] A. V. Manohar and M. B. Wise, Phys. Lett. B **636**, 107 (2006) [hep-ph/0601212].
- [20] A. Djouadi, Phys. Rept. **457**, 1 (2008) [hep-ph/0503172]; A. Djouadi, Phys. Rept. **459**, 1 (2008) [hep-ph/0503173].